

# From Strings to the LHC

Les Houches Lectures on String Phenomenology

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## Abstract

These notes are based on lectures given at the Les Houches Summer School in 2011, which was centered on the general topic *Theoretical Physics to face the challenge of LHC*. In these lectures I reviewed a number of topics in the field of String Phenomenology, focusing on orientifold/F-theory models yielding semi-realistic low-energy physics. The emphasis was on the extraction of the low-energy effective action and the possible test of specific models at LHC. These notes are a brief summary, appropriately updated, of some of the main topics covered in the lectures.

# 1 Branes and chirality

String Theory (ST) is the most serious candidate for a consistent theory of quantum gravity coupled to matter. In fact ST actually predicts the very existence of gravity, since a massless spin-2 particle, the graviton appears automatically in the spectrum of closed string theories. String Theory has also allowed us to improve our understanding of the origin of the blackhole degrees of freedom and also provides for explicit realizations of *holography* through the AdS/CFT correspondance. Remarkably, ST is not only a theory of quantum gravity but incorporates all the essential ingredients of the Standard Model (SM) of Particle Physics: gauge interactions, chiral fermions, Yukawa couplings... It is thus a strong candidate to provide us with a unified theory of all interactions, including the Standard Model (SM) and gravitation. In the last 25 years

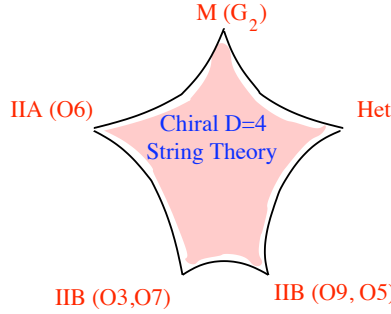


Figure 1: The five large classes of 4d chiral string compactifications.

enormous progress has been obtained in the understanding of the space of 4d string vacua [1]. From the point of view of unification, the main objective is to understand how the SM may be obtained as a low-energy limit of string theory. We would like to understand how the SM gauge group, the 3 quark/lepton generations, chirality, Yukawa couplings, CP-violation, neutrino masses, Higgs sector etc. may appear from an underlying string theory. The first step in that direction is learning which compactifications lead to a chiral spectrum of massless fermions at low-energies. There are essentially five large classes of such chiral 4d string vacua symbolized by the 5 vertices of the *pentagon* in fig.(1).

These include three large classes of Type II *orientifolds* (IIA with O6 orientifold planes, and IIB with O3/O7 or O9/O5 orientifold planes). In addition there are the well studied heterotic vacua in Calabi-Yau (CY) manifolds. Finally there are less studied (and difficult to handle) vacua obtained from the 11d M-theory compactified in manifolds of  $G_2$  holonomy. Different dualities connect these different corners so

the different classes of vacua should be considered as 5 different corners of a single underlying class of theories. It is impossible to overview all these different classes of theories so that we will concentrate on the case of the Type II orientifolds whose potential for the construction of realistic SM-like compactifications has been explored in the last 15 years.

The essential objects in chiral Type II orientifolds are  $Dp$ -branes, non-perturbative solitonic states of string theory which extend over  $(p+1)$  space+time dimensions. For our purposes  $Dp$ -branes may be considered as subspaces of the 10d space of Type II string theory in which open strings are allowed to start and end. They are charged under antisymmetric tensors of the Ramond-Ramond (RR) sector of Type II theory with  $(p+1)$  indices. Since in Type IIA(IIB) the massless RR tensors have an odd(even) number of indices, there are  $Dp$ -branes with  $p$  even(odd) for Type IIA(IIB) string theory. We will be interested in  $Dp$ -branes large enough to contain the standard Minkowski space inside so that the relevant  $Dp$ -branes will be  $D4, D6, D8$  in Type IIA and  $D3, D5, D7, D9$  in Type IIB. In compactified theories Gauss theorem will force the overall RR charges with respect to these antisymmetric fields to vanish. This leads to the so called *tadpole* cancellation conditions which turn out to also insure cancellation of gauge and gravitational anomalies in the theory.

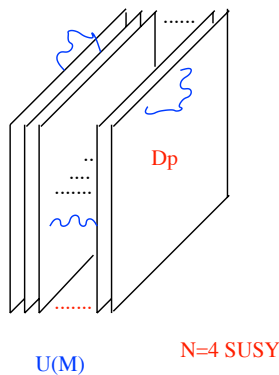


Figure 2: Open strings ending on a stack of  $M$  parallel  $Dp$ -branes give rise to a  $U(M)$ ,  $N=4$  gauge theory.

In the worldvolume of  $Dp$ -branes there live (are localized) gauge and charged matter degrees of freedom. In a single  $Dp$ -branes lives a  $U(1)$  gauge boson and  $M$  such branes located in the same place in transverse dimensions contain an enhanced  $U(M)$  gauge symmetry with  $N = 4$  SUSY in flat space. The corresponding spectrum is obviously

non-chiral and insufficient to yield realistic physics. In order to obtain chirality additional ingredients must be present. In the case of Type IIA models with the six extra dimensions compactified in a Calabi-Yau (CY) manifold, chiral fermions appear at the intersection of pairs of D6-branes, as we will describe later. In the case of Type IIB models chiral fermions may appear at the worldvolume of D7 or D9 branes in the presence of magnetic fluxes in the compact directions. Alternatively, chirality may appear if the geometry is singular, like e.g. the case of D3-branes on  $\mathbf{Z}_N$  orbifold singularities.

The other crucial ingredient in perturbative Type II models are Op *orientifolds*. These are geometrically analogous to Dp-branes with the crucial difference that they are not dynamical and do not contain any field degrees of freedom in their worldvolume. They are however charged under the RR antisymmetric fields and they have also *negative* tension compared to their Dp-branes counterparts. It is precisely these two properties which make useful the presence of orientifold planes, their negative tension and RR charges may be used to cancel the positive contribution of Dp-branes, allowing for the construction of Type II vacua with zero vacuum energy (Minkowski) and overall vanishing RR charges in a compact space.

Another important property of Type IIA and IIB vacua is Mirror symmetry. This is a symmetry which exchanges IIA and IIB compactifications by exchanging accordingly the underlying CY space by its mirror. For each CY manifold one can find a mirror manifold in which the Kahler and complex structure moduli are exchanged. In simple examples (like tori and orbifolds thereof) one can show that mirror symmetry is a particular example of *T-duality*. The action of T-duality in these toroidal/orbifold settings (to be discussed below) is non-trivial and exchanges Neumann and Dirichlet open string boundary conditions. An odd number of T-dualities along one-cycles exchanges Type IIA and IIB theories and the dimensionalities of Dp-branes changes accordingly. Thus e.g. 3 T-dualities on Type IIB D9-branes on  $T^6$  change them into D6 branes wrapping a 3-cycle in  $T^6$ .

The basic rules for D-brane model building are as follows [2]. One starts with Type II theory compactified on a CY (in some simple examples one may consider  $T^6$  tori or orbifolds). One then consider possible distributions of Dp-branes containing Minkowski space and preserving  $N = 1$  SUSY in 4d. The branes wrap subspaces (cycles) or are located at specific regions inside the CY. The configuration so far has positive energy and RR charges and is untenable if one wants to obtain Minkowski vacua. To achieve that, appropriate Op orientifold planes will be required both to cancel the positive vacuum energy and overall RR charges. This will require the construction of

a CY orientifold. Finally, the brane distribution is so chosen that the massless sector resembles as much as possible the SM or the MSSM. If the brane distribution respects the same  $N = 1$  SUSY in 4d the theory will be perturbatively stable.

In the above enterprise two approaches are possible:

- Global models. One insists in having a complete globally consistent CY compactification, with all RR tadpoles canceling.
- Local models. One considers local sets of lower dimensional Dp-branes ( $p \leq 7$ ) which are localized on some region of the CY and reproduce the SM or MSSM physics there. One does not care at this stage about global aspects of the compactification and assumes that eventually the configuration may be embedded inside a fully consistent global compact model.

The latter is often called the *bottom-up* approach [3], since one first constructs the local (bottom) model with the idea that eventually one may embed it in some global model. Note that this philosophy is not applicable to heterotic or Type I vacua since in those strings the SM fields live in the bulk six dimensions of the CY.

## 2 Type II orientifolds: intersections and magnetic fluxes

In Type IIA compactifications in principle we have D4, D6 and D8-branes, big enough to contain Minkowski space  $M_4$ . They can span  $M_4$  and wrap respectively 1-, 3- and 5-cycles in the CY. However, since CY manifolds do not have non-trivial 1- or 5-cycles, in IIA orientifolds only D6-branes are relevant for our purposes. It is easy to see that a pair of intersecting branes,  $D6_a$ ,  $D6_b$  give rise to chiral fermions at their intersection from open strings starting in one and ending on the other brane (see fig.(3)). The mass formula for the fields at an intersection in flat space is given (in bosonized formulation) by

$$M_{ab}^2 = N_{osc} + \frac{(r + r_\theta)^2}{2} - \frac{1}{2} + \sum_{i=1}^3 \frac{1}{2} |\theta_i| (1 - |\theta_i|) \quad (2.1)$$

where  $r_\theta = (\theta_1, \theta_2, \theta_3, 0)$  and  $r$  belongs to the  $SO(8)$  lattice ( $r_i = \mathbf{Z}, \mathbf{Z} + 1/2$  for NS, RR sectors respectively, with  $\sum_i r_i = \text{odd}$ ). The reader can check that the state  $r + r_\theta = (-\frac{1}{2} + \theta_1, -\frac{1}{2} + \theta_2, -\frac{1}{2} + \theta_3, +\frac{1}{2})$  is massless for any value of the angles, so there is always a massless fermion at the intersection. If there are N  $D6_a$  and M  $D6_b$  intersecting stacks of branes the fermion transforms in the bi-fundamental  $(N, \overline{M})$ .

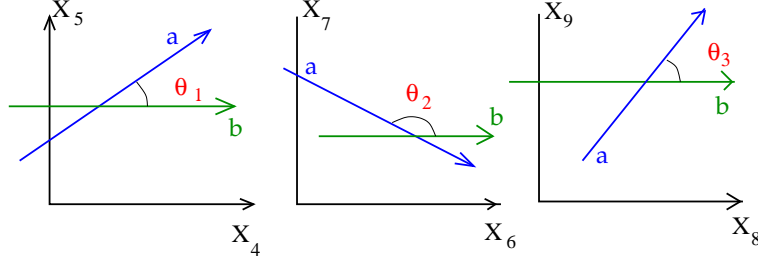


Figure 3: Open strings between  $D6_a$ - $D6_b$  branes intersecting at angles yield massless chiral fermions. Here  $X_5, \dots, X_9$  are local coordinates for the compact space.

There are also three scalars (e.g.  $r + r_\theta = (-1 + \theta_1, \theta_2, \theta_3, 0)$ ) with mass<sup>2</sup> which may be positive, zero or negative, depending on the values of the angles. Tachyons are avoided for large ranges of the intersecting angles. On the other hand for particular choices of the angles there is a massless scalar, the partner of the chiral fermion, signaling the presence of a  $N = 1$  SUSY, at least at the local level.

In order to construct 4d models one compactifies Type IIA string theory down to four dimensions on a CY manifold. The resulting theory has  $N = 2$  supersymmetry and is not yet suitable for realistic model building. One then constructs an orientifold by moding the theory by  $\Omega\mathcal{R}$  where  $\Omega$  is the worldsheet parity operation and  $\mathcal{R}$  is a  $\mathbf{Z}_2$  antiholomorphic involution on the CY with  $\mathcal{R}J = -J$  and  $\mathcal{R}\Omega_3 = \overline{\Omega}_3$  ( $J$  and  $\Omega_3$  are the Kahler 2-form and the holomorphic 3-form characteristic of CY manifolds). The resulting theory has now  $N = 1$  SUSY in 4d and the submanifolds left fixed under the  $\mathcal{R}$  operation are orientifold O6-planes carrying  $C_7$  RR antisymmetric field charge. To flesh out these process let us consider the simplified (yet phenomenologically interesting) case of a  $T^6$  orientifold compactification [4].

Consider Type IIA string theory compactified in a factorized torus  $T^6 = T^2 \times T^2 \times T^2$ . D6-branes are assumed to wrap  $\mathbf{M}_4$  and a 3-cycle which is the direct product of 3 1-cycles, one per  $T^2$  (see fig.4). These cycles are described by integers  $(n_a^i, m_a^i), i = 1, 2, 3$  indicating the number of times  $n_a^i(m_a^i)$  the  $D6_a$  brane wraps around the horizontal(vertical) directions. For each stack of  $N_a$   $D6_a$ -branes there is a  $U(N_a)$  gauge group. Furthermore at the intersection of two stacks of branes  $D6_a, D6_b$  the exchange of open strings gives rise to massless chiral fermions in bifundamental  $(\mathbf{N}_a, \overline{\mathbf{N}}_b)$  representations. Their multiplicity is given by their intersection number

$$I_{ab} = I_{ab}^1 \times I_{ab}^2 \times I_{ab}^3 = (n_a^1 m_b^1 - m_a^1 n_b^1)(n_a^2 m_b^2 - m_a^2 n_b^2)(n_a^3 m_b^3 - m_a^3 n_b^3). \quad (2.2)$$

which is e.g.  $2 \times 2 \times 1 = 4$  in the example of fig.(4). We now construct an orientifold

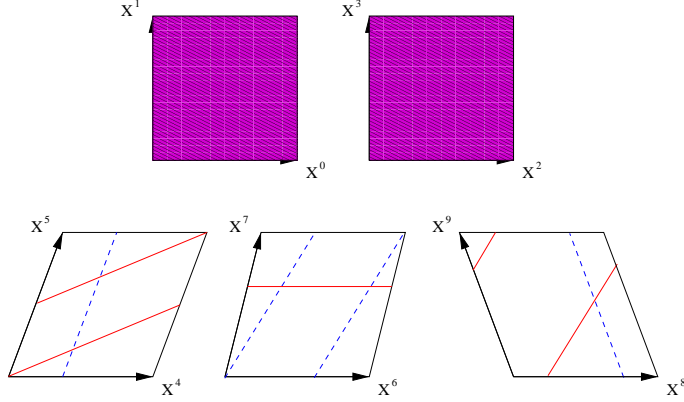


Figure 4: D6-branes wrap 1-cycles in each of the 3  $T^2$  and intersect at angles. Chiral bifundamental fermions are localized at the intersections.

by moding out the theory by the world-sheet operator  $\Omega(\tau, \sigma) = (\tau, -\sigma)$  acting on the world-sheet coordinates. Simultaneously we act with a reflection on the three coordinates  $R(X_i) = -X_i$ ,  $i = 5, 7, 9$ . This geometrical reflection leaves invariant the space defined by  $X_5 = X_7 = X_9 = 0$  in which the O6 orientifold lives. In addition the orientifold projection on invariant states may modify the gauge group of the branes if the latter wrap a 3-cycle which is left invariant by the orientifold. Depending on the details of the projection one may get  $Sp(N)$  or  $O(N)$  groups. On the other hand, if the 3-cycle wrapped by the D6-brane stack is not invariant, one must add in the background extra *mirror* D6\* branes sitting on the reflected 3-cycle with wrapping numbers  $(n^i, -m^i)$ . Then the configuration is also invariant but the gauge group  $U(N)$  remains.

It is easy to find choices of D6-branes with appropriate wrapping numbers  $(n^i, m^i)$  yielding a semirealistic chiral massless spectrum. Let us consider 3 stacks D6<sub>a</sub>, D6<sub>b</sub>, D6<sub>c</sub> of branes on rectangular  $T^2$  tori with multiplicities and wrapping numbers as in table 1 [5]. The D6<sub>b</sub> and D6<sub>c</sub> branes are assumed to be located at  $X_7 = 0$  and  $X_9 = 0$

$N_\alpha$	$(n_i^1, m_i^1)$	$(n_i^2, m_i^2)$	$(n_i^3, m_i^3)$
$N_a = 3 + 1$	(1, 0)	(3, 1)	(3, -1)
$N_b = 1$	(0, 1)	(1, 0)	(0, -1)
$N_c = 1$	(0, 1)	(0, -1)	(1, 0)

Table 1: Wrapping numbers of D6-branes in a MSSM-like configuration.

respectively so that the orientifold projection yields an  $Sp(1) \simeq SU(2)$  gauge group for both of them. On the other hand the 3 + 1 D6<sub>a</sub> branes in the table should be

supplemented by their mirrors with wrapping numbers flipped as  $(n_a^i, m_a^i) \rightarrow (n_a, -m_a^i)$ , and the gauge group is  $U(3) \times U(1)$ . The complete gauge group is then  $U(3) \times SU(2) \times SU(2) \times U(1)$  but one linear combination of the two  $U(1)$ 's is anomalous and becomes massive through a generalized Green-Schwarz mechanism. All in all, one obtains the gauge group of the minimal left-right symmetric extension of the MSSM. The reader may check using eq.(2.2) that one has three generations of quarks and leptons, with three right-handed neutrinos. Furthermore, if the branes  $D6_b$  and  $D6_c$  sit on top of each other in the first  $T^2$ , there is one minimal set of Higgs fields. Choosing  $R_x^2/R_y^2 = R_x^3/R_y^3$  for the radii in the second and third torus one can see that  $\theta_2 + \theta_3 = 0$  and there is one unbroken  $N = 1$  SUSY.

The above example is a *good local model* but it is globally inconsistent. The reason is that as it stands it gives rise to RR tadpoles, the overall charge with respect to the  $C^7$  RR forms does not vanish as it should in a compact space. It is easy to show that those conditions in this toroidal setting are

$$\sum_a N_a n_a^1 n_a^2 n_a^3 = 16 \ ; \ \sum_a N_a n_a^1 m_a^2 m_a^3 = 0 \ (+ \text{permutations}) \quad (2.3)$$

and plugging the wrapping numbers of the table one observes they are not obeyed. It is however easy to construct a  $\mathbf{Z}_2 \times \mathbf{Z}_2$  orbifold variation of this model with some additional D6-branes and orientifold planes which is supersymmetric and obeys the corresponding tadpole conditions [6].

This model is remarkably simple and its chiral sector gets quite close to a phenomenologically interesting model, the L-R extension of the MSSM. Still has the shortcoming that, like most toroidal/orbifold models, the massless spectrum includes additional adjoint chiral multiplets of the SM gauge group. The vev of these adjoints parametrize the freedom to translate in parallel the positions of the branes in any of these models. The latter is a characteristic of toroidal compactifications and is in general absent in more general CY orientifolds.

A second class of interesting Type II compactifications is Type IIB orientifolds. Now the internal orientifold geometric involution acts like  $RJ = J$ ,  $R\Omega_3 = -\Omega_3$ . In the toroidal setting they may be obtained as T-duals of Type IIA intersecting brane models. Indeed, upon an odd number of T-dualities along the 6 circles in the  $T^6$  a D6-brane may transform into a D9, D7, D5 or D3-brane, depending on the particular T-duality transformation. If the original D6 brane is rotated with respect to the orientifold plane the resulting IIB Dp-branes will in general contain a magnetic flux turned on in their worldvolume. Indeed, higher dimensional Type IIB branes in SUSY configurations (unlike the D6-branes in IIA) may contain magnetic flux backgrounds. They in turn



induce lower dimensional Dp-brane charge and also chirality. Let us consider [7] the case of  $N_a$  D9-branes wrapped  $m_a^i$  times on the  $i$ -th  $T^2$  and with  $n_a^i$  units of  $U(1)_a$  quantized magnetic flux:

$$m_a^i \frac{1}{2\pi} \int_{T^2} F_a^i = n_a^i . \quad (2.4)$$

Interestingly enough the  $(n_a^i, m_a^i)$  D6 wrapping numbers are mapped under T-duality into the magnetic integers defined above. In addition the relative angle  $\theta_{ab}^i$  of D6 <sub>$a$</sub> -D6 <sub>$b$</sub>  branes in the  $i$ -th torus is mapped into the difference

$$\theta_{ab}^i = \arctg(F_b^i) - \arctg(F_a^i) \quad , \quad F_a^i = \frac{n_a^i}{m_a^i R_{xi} R_{yi}} . \quad (2.5)$$

In the presence of a magnetic flux  $F$  in a IIB brane wrapping  $T^2$  the open string boundary conditions get modified as

$$\partial_\sigma X - F \partial_\tau Y = 0 \quad ; \quad \partial_\sigma Y + F \partial_\tau X = 0 . \quad (2.6)$$

In particular varying  $F$  one interpolates between Neumann and Dirichlet boundary conditions and e.g. at formally infinite flux they are purely Dirichlet. Thus adding fluxes on a higher dimensional brane induces RR charge corresponding to lower dimensional branes. For example, D9 branes with flux numbers  $(1,0)(n_a^2, m_a^2)(n_a^3, m_a^3)$  are equivalent to D7<sup>1</sup> branes which are localized on the first  $T^2$  and wrap the remaining  $T^2 \times T^2$ . On the other hand D9 branes with flux numbers  $(1,0)(1,0)(1,0)$  (formally infinite flux in the three  $T^2$ 's) is equivalent to D3-branes. Note in particular that the semirealistic model with intersecting D6-branes as in the table above are mapped into a set of three stacks of D7 <sub>$a$</sub> <sup>1</sup>, D7 <sub>$b$</sub> <sup>2</sup>, D7 <sub>$c$</sub> <sup>3</sup> which overlap pairwise on a  $T^2$ . Chirality arises in this Type IIB mirror from the mismatch between L- and R-handed fermions induced by the finite flux in the second and third tori.

This view of the orientifolds in terms of Type IIB D7-branes overlapping on dimension 2 spaces ( $T^2$  in the toroidal example) is particularly interesting because it admits a straightforward generalization to Type IIB CY orientifolds, at least in the large compact volume approximation in which Kaluza-Klein field theory techniques are available. On the contrary, the mirror class of models of Type IIA orientifolds with intersecting D6-branes is more difficult to generalize to curved CY spaces since the mathematical definition of BPS D6-branes in curved space (wrapping so called Special Lagrangian 3-cycles) is more difficult to analyze. A further argument to concentrate on Type IIB orientifolds with D7/D3 branes is that in the last 10 years we have learnt a great deal about how the addition of IIB closed string antisymmetric field fluxes can fix most or

all the moduli. The equivalent analysis for Type IIA or heterotic vacua is at present far less developed.

One generic problem in both IIA and IIB cases is the *top quark problem* in models with a unified gauge symmetry like  $SU(5)$ . The point is that in perturbative orientifolds the GUT symmetry is actually  $U(5)$  and the quantum numbers of a GUT generation are  $\bar{\mathbf{5}}_{-1} + \mathbf{10}_2$ , with Higgs multiplets  $\mathbf{5}_1 + \bar{\mathbf{5}}_{-1}$ . It is then clear that D-quark/lepton Yukawas are allowed by the  $U(1)$  symmetry but the U-quark couplings from  $\mathbf{10}_2\mathbf{10}_2\mathbf{5}_1$  are perturbatively forbidden. This  $U(1)$  symmetry is in fact anomalous and massive but still remains as a perturbative global symmetry in the effective action. Instanton effects may violate it but one expects the corresponding non-perturbative contributions to be small and be relevant at most only for the lightest generations, not the top quark. Thus insisting in unification of SM groups in perturbative orientifolds gives rise to a *top quark problem*.

### 3 Local F-theory GUT's

F-theory [8] may be considered as a geometric non-perturbative formulation of Type IIB orientifolds. From the model building point of view its interest is twofold: 1) It provides a solution to the *top quark problem* of perturbative Type II orientifolds with a GUT symmetry and 2) Moduli fixing induced by closed string antisymmetric fluxes is relatively well understood. In loose terms one could say that it combines advantages from the heterotic and Type IIB vacua.

An important massless field of 10d Type IIB string theory is the complexified dilaton field  $\tau = e^{-\phi} + iC_0$ . The dilaton  $\phi$  controls the perturbative loop expansion and  $C_0$  is a RR scalar. The 10d theory has a  $SL(2, \mathbf{Z})$  symmetry under which  $\tau$  as the modular parameter. The symmetry is generated by the transformations  $\tau \rightarrow 1/\tau$  and  $\tau \rightarrow \tau + i$  and is clearly non-perturbative (e.g. it exchanges strong and weak coupling by inverting the dilaton). F-theory provides a geometrization of this symmetry by adding two (auxiliary) extra dimensions with  $T^2$  geometry and identifying the complex structure of this  $T^2$  with the Type IIB  $\tau$  field. The resulting geometric construction is 12-dimensional and one obtains  $N = 1$  4d vacua by compactifying the theory on a complex 4-fold CY  $X_4$  which is an *elliptic fibration* over a 6-dimensional base  $B_3$ , i.e., locally one has  $X_4 \simeq T^2 \times B_3$ . The theory contains 7-branes which appear at points in the base  $B_3$  at which the fibration becomes singular, corresponding to 4-cycles wrapped by the 7-brane. As in the case of perturbative D7-branes, there is a gauge group

associated to these branes. However, unlike the perturbative case, the possible gauge groups include the exceptional ones  $E_6, E_7, E_8$ . This is an important property since, as we will see momentarily, allows for the existence of an  $SU(5)$  GUT symmetry with a large top Yukawa. A particularly interesting type of F-theory constructions are those

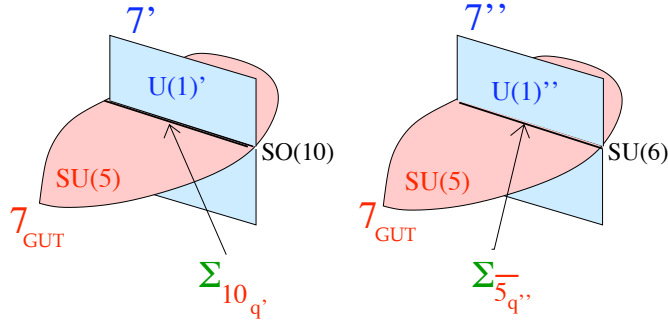


Figure 5: The  $SU(5)$  matter fields live at matter curves corresponding to the intersection of the bulk  $SU(5)$  brane with  $U(1)$  branes. At the matter curves the symmetry is enhanced to  $SU(6)$  and  $SO(10)$  respectively for the multiplets  $\mathbf{10}$  and  $\bar{\mathbf{5}}$ .

involving a GUT symmetry like  $SU(5)$ , termed F-theory GUT's [9]. This is motivated by the apparent unification of coupling constants in the MSSM. Such constructions are a non-perturbative generalization of the Type IIB models with intersecting (and magnetized) D7 branes that we discussed in previous section. There is a 7-brane wrapping a 4-cycle  $S$  inside  $B_3$  yielding an  $SU(5)$  gauge symmetry. As in the bottom-up approach mentioned above, one can decouple the local dynamics associated to the  $SU(5)$  brane from the global aspects of the  $B_3$  compact space. Chiral matter again appears at the intersection of pairs of 7-branes, *matter curves* in the F-theory language, corresponding to an enhanced degree of the singularity. These 7-branes are however non-perturbative and cannot simply be described in terms of perturbative open strings. A visual intuition of the arisal of matter fields in an  $SU(5)$  F-theory GUT is shown in fig.(5). At the matter curves the symmetry is locally enhanced to  $SU(6)$  or  $SO(10)$ . Recalling the adjoint branchings

$$\begin{aligned}
 SU(6) &\longrightarrow SU(5) \times U(1) & (3.1) \\
 \mathbf{35} &\longrightarrow \mathbf{24}_0 + \mathbf{1}_0 + [\mathbf{5}_1 + c.c.] \\
 SO(10) &\longrightarrow SU(5) \times U(1)'
 \end{aligned}$$

$$45 \longrightarrow 24_0 + 1_0 + [10_4 + c.c.] \quad (3.2)$$

one sees that in the matter curve associated to the 5-plet the symmetry is enhanced to  $SU(6)$  whereas in the one related to the 10-plets the symmetry is enhanced to  $SO(10)$ . Like in the perturbative magnetized IIB orientifolds, in order to get chiral fermions there must be in general non-vanishing fluxes along the  $U(1)$  and  $U(1)'$  symmetries. A third matter curve with an enhanced  $SU(6)'$  symmetry is also required to obtain Higgs 5-plets. Yukawa couplings appear at the intersection of the Higgs matter curve with the fermion matter curves, as illustrated in fig.(6). At the intersection point the symmetry is further enhanced to  $SO(12)$  in the case of the  $10 \times \bar{5} \times \bar{5}_H$  couplings and to  $E_6$  in the case of the U-quark couplings. One may now understand why there are U-quark Yukawa couplings in F-theory by looking at the branching of  $E_6$  adjoint into  $SU(5) \times U(1) \times U(1)'$ ,

$$\begin{aligned} E_6 &\longrightarrow SU(5) \times U(1) \times U(1)' \\ \mathbf{78} &\longrightarrow \text{Adjoint} + [(\mathbf{10}, -1, -3) + (\mathbf{10}, 4, 0) + (\mathbf{5}, -3, 3) + (\mathbf{1}, 5, 3) + h.c.] . \end{aligned} \quad (3.3)$$

We now see that one can form a  $10 \times 10 \times 5$  coupling which is indeed allowed by the  $U(1)$  symmetries. We will come back to the issue of Yukawa couplings in F-theory local GUT's in the next section.

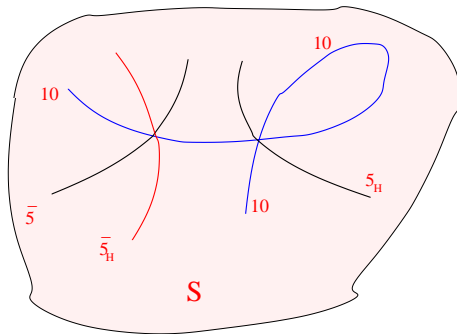


Figure 6: Matter curves generically intersect at points of further enhanced symmetry at which the  $10 \times \bar{5} \times \bar{5}_H$  and  $10 \times 10 \times \bar{5}_H$  Yukawa couplings localize.

To make the final contact with SM physics the  $SU(5)$  symmetry must be broken down to  $SU(3) \times SU(2) \times U(1)$ . In these constructions there are no massless adjoints to make that breaking and discrete Wilson lines are also not available. Still one can make such breaking by the addition of an additional flux  $F_Y$  along the hypercharge

direction of the  $SU(5)$ , which has the same symmetry breaking effect as an adjoint Higgs. Interestingly enough, this hypercharge flux may also be used to obtain doublet-triplet splitting of the Higgs multiplets  $\mathbf{5}_H + \bar{\mathbf{5}}_H$ .

## 4 The effective low energy action

To make contact with the low-energy physics we need to have information about the effective low-energy action remaining at scales well below the string scale. Here we will concentrate on the case of field theories with  $N = 1$  supersymmetry, which is assumed to be later broken at scales of order the Electro-Weak (EW) scale. In this case the action is determined by Kahler potential  $K$ , gauge kinetic functions  $f_a$  and the superpotential  $W$ , that we will discuss in turn. For definiteness we will concentrate also in the case of the effective action for Type IIB orientifolds, whose general features are also expected to apply for the F-theory case.

In the massless sector of a  $N = 1$  compactification there are charged fields from the open string sector (to be identified with the SM fields) and closed string fields giving rise to singlet chiral multiplets, the moduli. Among the latter there is the complex dilaton  $S = e^{-\phi} + iC_0$  which is just the dimensional reduction of the complex dilaton  $\tau$  mentioned above. In addition there are  $h_{11}$  Kahler moduli  $T^i$  and  $h_{21}^-$  complex structure moduli  $U^j$  (the minus means number of (2,1)-forms odd under the orientifold projection). The Kahler moduli parametrize the volume of the manifold and also of all the 4-cycles  $\Sigma_4^{(i)}$  of the specific CY. The complex structure fields  $U^j$  on the other hand parametrize the deformations of the CY manifolds and are associated to the 3-cycles  $\Sigma_3^{(j)}$  in the CY. Specifically one has [10] (in the simplest  $h_{11}^- = 0$  case)

$$T^i = e^{-\phi} Vol(\Sigma_4^{(i)}) + iC_4^{(i)} ; \quad U^j = \int_{\Sigma_3^{(j)}} \Omega_3 \quad (4.1)$$

Here  $C_4^{(i)}$  are 4d zero modes of the RR 4-form  $C_4$  on the 4-cycles. The  $N = 1$  supergravity Kahler potential associated to the moduli in Type IIB orientifold compactifications may be written as [10]

$$K_{IIB} = -\log(S + S^*) - 2\log(e^{-3\phi/2} Vol(CY)) - \log(-i \int \Omega_3 \wedge \bar{\Omega}_3), \quad (4.2)$$

where  $Vol(CY)$  is the volume of the CY manifold. In the toroidal case with rectangular  $T^6 = T^2 \times T^2 \times T^2$  the Kahler potential takes the simple form

$$K_{IIB} = -\log(S + S^*) - \sum_{i=1}^3 \log(U_i + U_i^*) - \sum_{i=1}^3 \log(T_i + T_i^*), \quad (4.3)$$

where  $T_i = e^{-\phi} R_x^j R_y^k R_x^k - iC_4$ , with  $i \neq j \neq k \neq i$  and  $U_i = R_y^i / R_x^i$ . This is the familiar no-scale structure which also appears in heterotic  $N = 1$  vacua.

Concerning the action for the charged matter fields  $\Phi_a$  on the 7-branes, the corresponding Kahler metrics, gauge kinetic functions and superpotential are themselves functions of the moduli. One can write for the general form of the supergravity Kahler potential an expression (to leading order in a matter field expansion)

$$K(M, M^*, \Phi_a, \Phi_a^*) = K_{IIB}(M, M^*) + \sum_{ab} K_{ab}(M, M^*) \Phi_a \Phi_b^* + \log |W(M) + W_Y(M, \Phi_a)|^2 \quad (4.4)$$

where  $M$  collectively denotes the moduli  $S, T^i, U^j$ ,  $W(M)$  is the superpotential of the moduli and  $W_Y(M, \Phi_a)$  the Yukawa coupling superpotential of the SM fields. We have already discussed the first term in eq.(4.4) and we will discuss the rest of the terms in what follows.

## 4.1 The Kahler metrics

Equation (4.4) includes the kinetic term for the matter fields which is controlled by the Kahler metrics  $K_{ab}$ , which is a function of the moduli. This dependence on the moduli is dictated by the geometric origin of the field. It has been computed at the classical level for some simple cases (mostly toroidal/orbifold orientifolds) either by dimensional reduction from the underlying 10d theory or using explicit string correlators. We are particularly interested in the Kahler metrics of fields living at intersecting 7-branes, since those are the ones which are associated to the MSSM fields in semi-realistic IIB or F-theory compactifications. In the case of type IIB toroidal/orbifold orientifolds the matter fields associated to a pair of intersecting  $D7^i$ - $D7^j$  branes has a metric (neglecting magnetic fluxes for the moment) [11]

$$K_{ab}^{ij} = \delta_{ab} \frac{1}{u_i^{1/2} u_j^{1/2} t_k^{1/2} s^{1/2}} \quad , \quad i \neq j \neq k \neq i \quad (4.5)$$

where  $t_i = (T_i + T_i^*)$ ,  $u_i = (U_i + U_i^*)$  and  $s = (S + S^*)$ . We thus see that the metrics of matter fields at untersections scale like  $K_{ab} \simeq t^{-1/2}$  with the Kahler moduli.

Toroidal orientifolds/orbifolds, however, are very special in some ways. We would rather like to see to what extent this type of Kahler metrics generalizes to more general IIB CY orientifolds. In particular  $D7$ -branes wrap 4-tori whose volumes are directly related to the overall volume of the compact manifold. One would rather like to obtain information about the Kahler metric when the 7-branes wrap a local 4-cycle whose volume is not directly connected to the overall volume of the CY. An example of

this is provided by the *swiss cheese* type of compactifications discussed in ref.[12]. In this more general setting one assumes that the SM fields are localized at D7-branes wrapping *small* cycles in a CY whose overall volume is controlled by a large modulus  $t_b$  (see fig.(7)) so the  $Vol[CY] = t_b^{3/2} - h(t_i)$ , where  $h$  is a homogeneous function of the *small* Kahler moduli  $t_i$  of degree 3/2. The simplest example of this is provided by the CY manifold  $\mathbf{P}^4_{[1,1,1,6,9]}$  which has only two Kahler moduli  $t_b, t$  with a Kahler potential of the form

$$K_{IIB} = -2\log(t_b^{3/2} - t^{3/2}) . \quad (4.6)$$

Here we will assume  $t_b \gg t$  and take both large so that the supergravity approximation

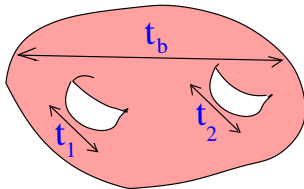


Figure 7: CY manifold with a *swiss cheese* structure.

is still valid. In the F-theory context the analogue of these moduli  $t, t_b$  would correspond to the size of the 4-fold  $S$  and the 6-fold  $B_3$  respectively. Focusing only in the Kahler moduli dependence of the metrics, one can write a large volume ansatz for the Kahler metrics of charged matter fields at the intersections [13]

$$K_\alpha = \frac{t^{(1-\xi_\alpha)}}{t_b}, \quad (4.7)$$

with  $\xi_\alpha$  to be fixed. One can compute  $\xi_\alpha$  by studying the behavior with respect to a scaling of  $t$  in the effective action. In particular in  $N = 1$  supergravity the physical (i.e. with normalized kinetic terms) Yukawa coupling  $\hat{Y}_{\alpha\beta\gamma}$  among three chiral fields are related to the holomorphic Yukawa coupling  $Y_{\alpha\beta\gamma}^{(0)}$  by

$$\hat{Y}_{\alpha\beta\gamma} = e^{K/2} \frac{Y_{\alpha\beta\gamma}^{(0)}}{(K_\alpha K_\beta K_\gamma)^{1/2}} . \quad (4.8)$$

On the other hand it is well known that the perturbative holomorphic Yukawa couplings in Type IIB string theory are independent of Kahler moduli. Then using eqs. (4.6),(4.7) one finds a scaling of the physical Yukawa

$$\hat{Y}_{\alpha\beta\gamma} \simeq t^{(\xi_\alpha + \xi_\beta + \xi_\gamma - 3)/2} . \quad (4.9)$$

The dependence on  $t_b$  drops at leading order in  $t/t_b$ , as expected for a model whose physics is essentially localized on the 4-cycle parametrized by  $t$ . On the other hand one

can alternatively compute the scaling behavior of the physical Yukawa in terms of its computation as an overlap integral of the respective wave functions in  $S$  (see below) so that

$$\hat{Y}_{\alpha\beta\gamma} \simeq \int \Psi_\alpha \Psi_\beta \Psi_\gamma \quad , \quad \int |\Psi_\alpha|^2 = 1. \quad (4.10)$$

For fields localized at intersecting branes the normalization integrals above are essentially 2-dimensional so that the wave functions should scale like  $t^{-1/4}$ . On the other hand the overlap integral for the Yukawa is essentially point-like so that it scales like  $\hat{Y} \simeq t^{-3/4}$ . Comparing this to eq.(4.9) one concludes that all  $\xi_\alpha = 1/2$  and hence the matter metrics of fields at intersecting branes have a metric with a local kahler modulus dependence of the form

$$K_\alpha = \frac{t^{1/2}}{t_b} . \quad (4.11)$$

Note that setting  $t_b \simeq t$  reproduces the Kahler modulus dependence  $t^{-1/2}$  of toroidal models eq.(4.5).

## 4.2 The gauge kinetic function

The gauge kinetic function for the gauge group living on the D7-worldvolume may be extracted by expanding the Dirac-Born-Infeld (DBI) action of the D7-brane to second order in the gauge field strength  $F_a$ . One obtains the general expression [1]

$$f_a^{\text{D7}} = \frac{(\alpha')^{-2}}{(2\pi)^5} \left( e^{-\phi} \int_{\Sigma_4^a} \text{Re} \left( e^{J+i2\pi\alpha' F_a} \right) + i \int_{\Sigma_4^a} \sum_k C_{2k} e^{2\pi\alpha' F_a} \right), \quad (4.12)$$

where  $J$  is the Kahler 2-form and the second piece performs a formal sum over all RR  $C_{2k}$  forms contributing to the integral. Upon expanding the exponential, the first term produces the volume of 4-fold  $\Sigma_4^a$  wrapped by the D7 and the second term is proportional to  $C_4$ . Taking into account eq.(4.1) one sees that the gauge kinetic function is proportional to the Kahler modulus  $T_a$ , i.e.

$$f_a = T_a . \quad (4.13)$$

The subsequent terms in the expansion describe contributions from the worldvolume gauge magnetic flux which will be subleading for large volume  $t_a = \text{Re}T_a$ . In particular, since  $\int_{\Sigma_2} F_a = n$ , with  $n = \text{integer}$  for quantized gauge fluxes, one can estimate the flux density as  $F_a \simeq n(\text{Re}S/\text{Re}T_a)^{1/2}$ . The leading flux correction to the gauge kinetic function has then the form

$$\text{Re}f_a \simeq \text{Re}T_a(1 + |F_a|^2) \simeq \text{Re}T_a(1 + n^2 \text{Re}(S)/\text{Re}(T_a)) \quad (4.14)$$

which indeed will be subleading in the large  $t_a$  limit.



### 4.3 The superpotential

As indicated in eq.(4.4) there will be a superpotential for the moduli  $W(M)$  and a second superpotential  $W(M, \Phi_\alpha)$  involving (moduli dependent) Yukawa couplings. In the absence of closed string antisymmetric fluxes the perturbative superpotential of the moduli vanishes,  $W_{pert}(M) = 0$ . However in the presence of Type IIB NS(RR) 3-form fluxes  $H_3(F_3)$  there is an induced effective superpotential involving the complex dilaton  $S$  and the complex structure moduli  $U^j$  given by [14]

$$W_{flux}(S, U^j) = \int_{CY} (F_3 - iSH_3) \wedge \Omega_3 \quad (4.15)$$

where  $\Omega_3$  is the CY holomorphic 3-form which may be expanded in terms of the complex structure moduli  $U^j$  in the CY. This superpotential depends only on  $S$  and the  $U^j$  and its minimization can give rise easily to the fixing of all these moduli. Furthermore since generic CY manifolds have of order of a hundred  $U^j$  fields or more, and the fluxes  $F_3, H_3$  have also a range of possible quantized values, at the minimum there may be accidental cancelations such that there is a very tiny value for  $\langle W_{flux} \rangle = W_0$ . Such tiny values would be needed to understand the smallness of the SUSY breaking scale compared to a large string scale  $M_s$  not much below the Planck scale.

The above fluxes are unable to fix the values of the Kahler moduli. However in specific compactifications there are non-perturbative effects which induce superpotential terms involving the Kahler moduli. Examples of such non-perturbative effects are instanton effects induced by euclidean  $D3$  instantons and gaugino condensation on  $D7$ -branes wrapping appropriate 4-cycles in the CY. Such effects have typically an exponentially suppressed behavior of the form  $W_{np} \simeq \sum_a \exp(-B_a T_a)$  for some constants  $B_a$ . These effects combined with those induced by fluxes  $W_{flux}$  have the potential to fix all the moduli of specific CY orientifold compactifications [15, 12]. Although a detailed example with all the required properties, including a realistic model and non-vanishing (but very small) cosmological constant, is still lacking, it seems very likely that those ingredients have the potential to fix all moduli.

Of more direct phenomenological interest are the Yukawa couplings involving SM quarks and leptons to Higgs scalars. As we already mentioned, Yukawa couplings among  $Dp$  brane matter fields in Type IIB compactification arise from overlap integral of the wave function in extra dimensions of the three participant fields. Consider the case of  $D9$ -branes to simplify the discussion (recall the case of  $D7$ -branes may be described in terms of  $D9$ -branes with appropriate fluxes). Suppose we have initially a  $D = 10$  Type IIB orientifold with  $D9$  branes and a gauge group  $U(n)$ . In the field

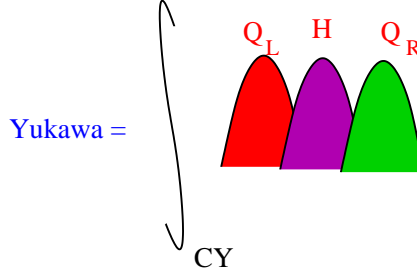


Figure 8: Pictorial representation of the computation of Yukawa coupling constants as overlap integrals of zero modes

theory limit our action will be 10d super-Yang-Mills,

$$L = -\frac{1}{4}Tr\left(F^{MN}F_{MN}\right) + \frac{i}{2}Tr\left(\bar{\Psi}\Gamma^M D_M\Psi\right). \quad (4.16)$$

We then compactify the theory on some CY manifold and turn on magnetic fluxes which may break the gauge group to a SM-like gauge group. The 10d fields can then be expanded a la Kaluza-Klein (KK)

$$\Psi(x^\mu, y^m) = \sum_k \chi_{(k)}(x^\mu) \otimes \psi_{(k)}(y^m), A_n(x^\mu, y^m) = \sum_k \varphi_{(k)}(x^\mu) \otimes \phi_{(k),n}(y^m) \quad (4.17)$$

where  $x^\mu$  and  $y^m$  are 4d and internal coordinates respectively. The 4d massless spectrum may be chiral and  $N = 1$  supersymmetric with judicious choice of magnetic fluxes. The 4d Yukawa coupling between these matter fields arise from KK reduction of the cubic coupling  $A \times \Psi \times \Psi$  from the 10d Lagrangian in eq.(4.16). As illustrated in fig.(8) the Yukawa coupling coefficients are obtained from the overlap integrals

$$Y_{ijk} = \frac{g}{2} \int_{CY} \psi_i^{\alpha\dagger} \Gamma^m \psi_j^\beta \phi_{km}^\gamma f_{\alpha\beta\gamma}, \quad (4.18)$$

where  $g$  is the 10d gauge coupling,  $\alpha, \beta, \gamma$  are  $U(n)$  gauge indices and  $f_{\alpha\beta\gamma}$  are  $U(n)$  structure constants; also  $\psi, \phi$  are fermionic and bosonic zero modes respectively, and  $i, j, k$  label the different zero modes in a given charge sector, i.e. the families in semi-realistic models. The Yukawa couplings are thus obtained as overlap integrals of the three zero mode wave functions in the CY.

In order to compute the Yukawa coupling constants we thus need to know the explicit form of the wave functions on compact dimensions of the involved matter fields, quarks, leptons and Higgs multiplets in a realistic model. However such wave functions are only accessible to explicit computation for simple models like toroidal compactifications or orbifolds thereof. Indeed this computation has been worked out for

general toroidal/orbifold models [16]. The wave functions turn out to be proportional to Jacobi  $\theta$ -functions with a Gaussian profile and the holomorphic Yukawa couplings turn out to be also proportional to products of Jacobi  $\theta$ -functions (one per  $T^2$  factor) with dependence only on the complex structure moduli  $U^j$  and the open string moduli (Wilson line degrees of freedom). As an example, the semirealistic model in table (1) has holomorphic Yukawa couplings with the structure

$$\begin{aligned} Y_{ij}^U &\sim \vartheta \begin{bmatrix} \frac{i}{3} \\ 0 \end{bmatrix} (3J^{(2)}) \times \vartheta \begin{bmatrix} \frac{j}{3} \\ 0 \end{bmatrix} (3J^{(3)}), \\ Y_{ij^*}^D &\sim \vartheta \begin{bmatrix} \frac{i}{3} \\ 0 \end{bmatrix} (3J^{(2)}) \times \vartheta \begin{bmatrix} \frac{j^*}{3} \\ 0 \end{bmatrix} (3J^{(3)}), \end{aligned} \quad (4.19)$$

where  $\vartheta$  is a Jacobi  $\theta$ -function,  $J^{(i)}$  are the Kahler forms of the  $i$ -th torus and  $i, j, j^*$  are family indices for the  $Q, U$  and  $D$  SM chiral multiplets respectively. These expressions yield proportional expressions for  $U$ - and  $D$ -quark Yukawa couplings but they differ if one takes into account the generic possibility (in tori) of Wilson line backgrounds along the  $T^6$  circles. However, due to the factorized structure of the family dependence, only one quark/lepton generation gets a mass. The corresponding Yukawa coupling is of order of the gauge coupling constant. This may be considered as a good first approximation to the observed quark/lepton mass spectrum and one expects further corrections to give rise to the Yukawa couplings of the lighter generations.

The computation of Yukawa couplings in general curved CY manifolds is more difficult, although it becomes more tractable within the context of the *bottom-up* approach mentioned above. The idea is that in models in which the SM fields are localized in brane intersections, Yukawa couplings appear at points in the CY in which three such intersections (corresponding to SM and Higgs fields) meet. We already saw that in the F-theory context in fig.5. Thus e.g. the Yukawa coupling  $\mathbf{10} \times \bar{\mathbf{5}} \times \bar{\mathbf{5}}_{\mathbf{H}}$  in a  $SU(5)$  F-theory GUT is localized at a point of triple intersection of the three matter curves. The Yukawa coupling has now the schematic form  $\int_S \psi_i \psi_j \phi_H$  in which the integral, extended over the 4-fold  $S$ , is dominated by the intersection region. In such a situation, to compute the Yukawa coupling we only need to know the wave functions *in the neighborhood* of the intersection point [17]. Those local wave functions may be obtained by solving the Dirac and K-G equations at the local level. Interestingly, one again finds that only one (the third) generation gets a non-vanishing Yukawa coupling, which is also of order the gauge coupling constant. It has been found however that instanton corrections induced by distant 7-branes wrapping other 4-cycles in compact space in general induce the required Yukawa couplings for the lighter generations [18, 19].

Instanton effects do not only give rise to superpotentials for the Kahler moduli and induce the Yukawa couplings of the lighter generations. They may also give rise to interesting terms in the SM superpotential which are forbidden in perturbation theory. In particular, in brane models of SM physics there are typically extra  $U(1)$  gauged symmetries beyond those of the SM. A classical example is  $U(1)_{B-L}$  which often appears gauged in many string constructions including right-handed neutrinos. This symmetry is anomaly-free but there are often in addition  $U(1)$ 's with triangle anomalies which are cancelled by the 4d version of the Green-Schwarz mechanism. All anomalous  $U(1)$ 's become massive by combining with the imaginary part of the Kahler (complex structure) moduli in Type IIB (IIA) orientifolds. But, in addition, anomaly-free gauge symmetries like  $U(1)_{B-L}$  may also become massive in this way. This happens due to the fact that e.g. in Type IIB some  $\text{Im}T_i$  transform under the corresponding gauge  $U(1)_a$  symmetries as  $\text{Im}T_i \rightarrow \text{Im}T_i + q_i^a \Lambda_a$ , with  $\Lambda_a$  the gauge parameter. This has interesting consequences for instanton physics [20]. In a Type IIB orientifold some instanton configurations corresponds to Euclidean D3-branes wrapping the compact dimensions (so that they are localized in Minkowski, as instantons should). If they intersect the D7-branes where the SM fields live, there appear charged zero modes (from open string exchange) contributing to instanton induced transitions. This is why this class of stringy instantons are often called *charged* instantons. In particular if the D3-brane wraps a 4-cycle with Kahler modulus  $M$  (some linear combination of the  $T_i$ 's), non-perturbative operators of the general form

$$e^{-M} \Phi_{q_1} \dots \Phi_{q_n} \quad , \quad \sum_i q_i \neq 0 \quad (4.20)$$

may appear [20]. These operators are gauge invariant because the sum of the charges of the  $\Phi$  chiral fields is compensated by the shift on  $\text{Im}M$  induced by the gauge transformation. An example of this is the generation of right-handed neutrino masses in MSSM-like orientifolds with a massive  $U(1)_{B-L}$  induced by a G-S mechanism. In this case the operator has the form  $e^{-M} \nu_R \nu_R$  and the non-invariance of the bilinear under  $U(1)_{B-L}$  is compensated by a shift of the  $\text{Im}M$ . The mass is of order  $e^{-\text{Re}M} M_s$  which may be on the right phenomenological ballpark  $10^{12} - 10^{14}$  GeV for  $M_s \simeq 10^{16}$  GeV and  $\text{Re}M \simeq 100$ . This type of *charged* instantons could also be important for the generation of other phenomenologically relevant terms like e.g. the MSSM  $\mu$ -term.

## 5 String model building and the LHC

With the LHC in operation an important issue is trying to make contact between an underlying string theory and experimental data. Of course it would be really exciting if the string scale  $M_s$  was within reach of the LHC. We could perhaps observe some string or KK excitation as resonances in LHC data. On the other a large string scale  $M_s \simeq 10^{16}$  GeV seems to be favored if one sticks to a SUSY version of the SM such as the MSSM, in which gauge couplings nicely unify at a scale of order  $10^{16}$  GeV. So it is important to see whether specific classes of string compactifications may lead to low energy predictions for SUSY breaking parameters.

We have seen that in certain large classes of Type II models there is information about the structure of the low-energy effective action. In particular in Type IIB orientifolds (or their F-theory extension) with SM fields localized at intersecting D7-branes (or matter curves in F-theory GUT's) one can compute the dependence on the local Kahler modulus of the gauge kinetic function (eq.(4.13)) and also of the Kahler metric (eq.(4.11)). If a MSSM-like model is constructed on such a setting, one can obtain specific expressions for SUSY breaking soft terms assuming Kahler moduli dominance in SUSY breaking, i.e., non-vanishing auxiliary fields  $F_t \neq 0$ . This is a reasonable assumption within Type IIB/F-theory since in Type IIB orientifolds such non-vanishing auxiliary fields correspond to the presence of non-vanishing antisymmetric RR and NS imaginary self-dual  $(0, 3)$  fluxes [21], which are known to solve the classical equations of motion [22]. As we mentioned above, such closed string fluxes are generically present in compactifications with fixed moduli. Using standard  $N = 1$  supergravity formulae and the above information on the effective action one obtains soft terms with the CMSSM structure but with the additional relationships [23],

$$M = \sqrt{2}m = -(2/3)A = -B. \quad (5.1)$$

where  $M$  is the universal gaugino mass,  $m$  the universal scalar mass,  $A$  the trilinear scalar parameter and  $B$  the Higgs bilinear parameter. Here one assumes the presence of an explicit  $\mu$ -term in the low energy Lagrangian so that altogether there are only 2 free parameters,  $M$  and  $\mu$ . The universality of soft terms may be understood if an underlying GUT structure exists as in F-theory GUT's. As we have mentioned, magnetic flux backgrounds are generically present on the worldvolume of the underlying 7-branes in order to get a chiral spectrum. In the presence of magnetic fluxes the gauge kinetic functions (see eq.(4.14)) and the Kahler metrics may get small corrections to

eqs.(4.13),(4.11), i.e.

$$f = T(1 + a\frac{S}{T}) ; K_\alpha = \frac{t^{1/2}}{t_b}(1 + \frac{c_\alpha}{t^{1/2}}) , \quad (5.2)$$

where  $a$  and  $c_\alpha$  are constants and  $S$  is the the complex dilaton field. These corrections are suppressed in the large  $t$  limit, corresponding to the physical weak coupling. In this limit one may also neglect the correction to  $f$  compared to that coming from  $K_\alpha$ . One then finds corrected soft terms of the form

$$m_{\tilde{f}}^2 = \frac{1}{2}|M|^2(1 - \frac{3}{2}\rho_f) , \quad (5.3)$$

$$m_H^2 = \frac{1}{2}|M|^2(1 - \frac{3}{2}\rho_H) , \quad (5.4)$$

$$A = -\frac{1}{2}M(3 - \rho_H - 2\rho_f) , \quad (5.5)$$

$$B = -M(1 - \rho_H) , \quad (5.6)$$

where  $\rho_\alpha = c_\alpha/t^{1/2}$  and the subindices  $f, H$  refer to fluxes through the fermion matter curves or the Higgs curve. Note that as an order of magnitude one numerically expects  $\rho_{H,f} \simeq 1/t^{1/2} \simeq \alpha_{GUT}^{1/2} \simeq 0.2$ . The above soft terms apply at the string/unification scale  $M_s \simeq 10^{16}$  GeV. In order to get the low energy physics around the EW scale one has to run down the soft parameters according to the renormalization group equations (RGE). Then one has to check that the boundary conditions are consistent with radiative EW symmetry breaking (REWSB) and with present low-energy phenomenological constraints. One may in addition impose that the lightest neutralino is stable and provides for the dark matter in the universe. The resulting scheme is extremely constrained [24]. In particular, setting the fermion flux correction to zero for simplicity, one has a theory with three free parameters ( $M, \mu$  and  $\rho_H$ ) and two constraints (REWSB and dark matter), or equivalently, lines in the planes of any pair of parameters or SUSY masses. As an example fig.(9) shows the normalized mass difference  $(m_{\tilde{\tau}_1} - m_{\chi_1^0})/m_{\tilde{\tau}_1}$  as a function of the lightest Higgs mass  $m_h$  [24]. Dots correspond to points fulfilling the central value in the result from WMAP for the neutralino relic density and dotted lines denote the upper and lower limits after including the  $2\sigma$  uncertainty. The dot-dashed line represents points with a critical matter density  $\Omega_{matter} = 1$ . The vertical line corresponds to the  $2\sigma$  limit on  $\text{BR}(b \rightarrow s\gamma)$  and the upper bound on  $\text{BR}(B_s \rightarrow \mu^+\mu^-)$  from Ref. [25] and the recent LHCb result [26]. The gray area indicates the points compatible with the latter constraint when the  $2\sigma$  error associated to the SM prediction is included. As is obvious from the figure, the dark matter condition is fulfilled thanks to a stau-neutralino coannihilation mechanism. Interestingly enough, the recent constraint on  $\text{BR}(B_s \rightarrow \mu^+\mu^-)$  from LHCb forces the Higgs mass to a region around 125

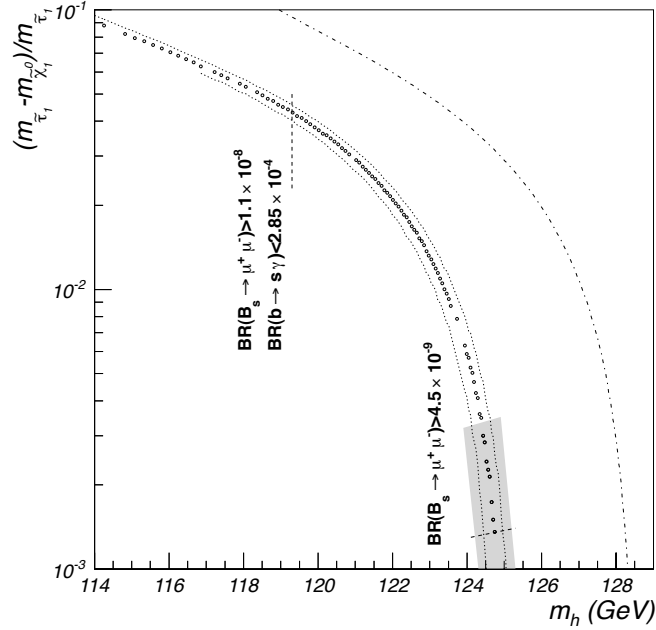


Figure 9: The normalized mass difference  $(m_{\tilde{\tau}_1} - m_{\chi_1^0})/m_{\tilde{\tau}_1}$  as a function of the lightest Higgs mass  $m_h$  in the modulus dominance scheme. Appropriate REWSB, neutralino dark matter and  $\text{BR}(B_s \rightarrow \mu^+ \mu^-)$  limits are only consistent for a Higgs mass in the 125 GeV region. (from ref.[24]).

GeV, consistent with the hints of a Higgs particle in that range as measured at CMS and ATLAS. Fixing the mass of any SUSY particle fixes the rest of the spectrum. In particular, with a lightest Higgs mass around 125 GeV gluinos have a mass around 3 TeV, the 1-st,2-nd generation squarks around 2.7-2.8 TeV and the lightest stop around 2 TeV. The lightest slepton is a stau with mass around 600 GeV, almost degenerate with the lightest neutralinos. The existence of gluino and squarks of these masses can be tested at LHC running at 14 TeV and  $30 \text{ fb}^{-1}$  integrated luminosity.

It is remarkable that a lightest MSSM Higgs mass as heavy as 125 GeV is possible in this scheme. In most SUSY schemes (including minimal gauge and anomaly mediation models and the CMSSM with not superheavy squarks) the lightest Higgs mass is typically around 115 GeV or so [27]. In this scheme a relatively heavy Higgs appears because the soft terms in eq.(5.6) predict a large  $A$ -parameter with  $A \simeq -2m_t$ , giving rise to a large stop mixing parameter and hence a big one-loop correction to the Higgs mass. In addition the dark matter and REWSB conditions require a large  $\tan\beta \simeq 40$ ,

pushing the tree level Higgs mass to its maximum value. This large  $\tan\beta$  and stop mixing parameters imply that, as it stands this simple scheme may be soon ruled out if LHCb finds no deviation from the SM value for  $\text{BR}(B_s \rightarrow \mu^+ \mu^-)$ , which may happen soon. On the other hand a NMSSM version of the same model, also viable in Type IIB/F-theory schemes, would remain consistent as would also R-parity violation, since it would avoid the dark matter over-abundance problem. This shows how the LHC results may provide important constraints on the possible compactifications and SUSY-breaking schemes within string theory, see e.g. ref.([28]) for other string derived approaches.

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